Electromagnetic extraction of energy from black hole-neutron star binaries

Sean T. McWilliams[1,2,3] and Janna Levin[1,3]

[1] *Institute for Strings, Cosmology and Astroparticle Physics (ISCAP), Columbia University, New York, NY 10027*

[2] *Physics Department, Princeton University, Princeton, NJ 08544 and*

[3] *Department of Physics and Astronomy, Barnard College, New York, NY 10027*[*]

(Dated: January 12, 2011)

The coalescence of black hole-neutron star binaries is expected to be a principal source of gravitational waves for the next generation of detectors, Advanced LIGO and Advanced Virgo. Ideally, these and other gravitational wave sources would have a distinct electromagnetic counterpart, as significantly more information could be gained through two separate channels. In addition, since these detectors will probe distances with non-negligible redshift, a coincident observation of an electromagnetic counterpart to a gravitational wave signal would facilitate a novel measurement of dark energy [1]. For black hole masses not much larger than the neutron star mass, the tidal disruption and subsequent accretion of the neutron star by the black hole provides one avenue for generating an electromagnetic counterpart [2]. However, in this work, we demonstrate that, for all black hole-neutron star binaries observable by Advanced LIGO/Virgo, the interaction of the black hole with the magnetic field of the neutron star will drive a Poynting flux. This Poynting flux generates synchrotron/curvature radiation as the electron-positron plasma in the neutron star magnetosphere is accelerated, and thermal radiation as the plasma is focused onto the neutron star magnetic poles, creating a "hot spot" on the neutron star surface. This novel effect will generate copious luminosity, comparable to supernovae and active galactic nuclei, so that black hole-neutron star coalescences detectable with gravitational waves by Advanced LIGO/Virgo could also potentially be detectable electromagnetically.

*Electronic address: sean@astro.columbia.edu

When a neutron star is tidally disrupted by a stellar-mass black hole, the system will release substantial amounts of energy as the remnant material is accreted. Along with neutron star-neutron star binary mergers, these events are a leading candidate for driving short γ-ray bursts (GRBs), which persist for less than ~ 2 seconds, yet can release more energy in that time than most galaxies will emit in an entire year. Apart from being the most luminous occurence in the Universe, these short GRBs are potentialy valuable electromagnetic counterparts to gravitational wave observations. However, unless the black hole in these systems is of sufficiently low mass and sufficiently high spin, the neutron star will simply be swallowed whole. Nonetheless, even when a neutron star survives a merger with a black hole intact, there is still a promising source for a bright electromagnetic event that has not previously been discussed. In this work, we will show that when a black hole and neutron star are close enough that the black hole orbits within the neutron-star magnetosphere, and the magnetic field threads the black-hole horizon, a circuit is established with the neutron star as an external resistor, the magnetic field lines as wires, the magnetospheric electron-positron plasma providing current, and the black hole acting as a battery. The result, as we will show, is a γ-ray burst of very short duration which should nonetheless be detectable at cosmological distances with sufficient timing resolution. Since this process does not depend on the disruption of the neutron star to operate, these events may prove more common than short GRBs generated by tidal disruption and accretion.

In the well-known Blandford-Znajek scenario [3] for driving electromagnetic jets in active galactic nuclei, the spinning supermassive black hole at the center of a galaxy is embedded in a magnetic field. The field is assumed to be anchored in an accretion disk, which is further assumed to be an excellent conductor, so that the field lines are affixed to the disk. The spin of the black hole then twists the magnetic field lines, driving a Poynting flux at the expense of its own spin. Recently, it has been shown [4–6] that this picture could also be applied to a black hole-black hole binary, orbiting through a fixed background magnetic field, where the orbital energy drives the Poynting flux. However, it is unclear how magnetic fields as strong as is necessary ($B \approx 10^4$ G) can be expected to surround an orbiting binary, since the binary will exert a gravitational torque that counteracts the viscous inflow of the accretion disk [7], thereby inhibiting the deposition of a magnetic field on the black hole. It is conceivable that plasma from the disk that does manage to accrete onto an orbiting binary will tow its magnetic field along, but it seems less likely that this will cause the magnetic

field around the orbiting binary to be comparable in strength to the field found in the bulk of the accretion disk.

A different but related scenario, which will be the focus of this work, is the Poynting flux generated by the inspiral of a black hole-neutron star binary. Much attention has been given to the possibility that the neutron star will disrupt, and the remnant gas will be accreted (see [2] for the theoretical description, and [8–10] for results from recent simulations). However, since neutron stars can possess surface magnetic fields of $\mathcal{O}(10^{12}$ G) (and $\sim 10^{15}$ G in the case of magnetars), it is conceivable that they could provide a mechanism for driving a substantial Poynting flux. Since the neutron star core that anchors the magnetic field is superconducting, the field lines will be affixed to the neutron star, and the relative motion and/or rotation of the neutron star and the orbiting black hole will generate a Poynting flux. Furthermore, since we know that neutron stars can possess such large magnetic fields, the inspiral and eventual merger of black hole-neutron star binaries requires no *ad hoc* mechanism for transporting a sufficiently strong magnetic field to the vicinity of the black hole. It is worth noting that the magnetic fields of neutron stars may decay on shorter timescales ($\sim 10^7$ years [11]) than the typical inspiral timescale (10^9 years [12]), so that the neutron star might have a smaller magnetic field of $\mathcal{O}(10^8$ G) at merger, which would significantly mitigate the effect we describe. However, the decay timescale quoted above assumes there is no replenishment of the magnetic field, for example through dynamo action, and that the field is entirely the result of amplification of the progenitor star's field during the collapse to a neutron star. Furthermore, it is possible that the tidal deformation of the neutron star during the late inspiral could result in significant magnetic field amplification [8]. Even in the most pessimistic scenario, provided the neutron star magnetic field exceeds 10^8 G at merger, the plasma will still be channeled to the magnetic poles by the magnetic field [13], and the mechanism we describe will still operate, albeit at lower magnitudes.

To simplify the calculation, we employ the membrane paradigm [14], wherein the black hole event horizon is reinterpreted as a two-dimensional conducting membrane, and the standard Maxwell equations in three dimensions govern the dynamics of electromagnetic fields evolving along a sequence of hypersurfaces with minimal modification. Using this approach, the Blandford-Znajek effect for a single spinning black hole immersed in plasma can be calculated using simple circuit equations. In that setting, the black hole spin drives a current along the magnetic field lines by Faraday induction. For the case we are addressing, Fara-

day induction results from a combination of the orbital motion and spin of the black hole through the dipole magnetic field of the neutron star. As the inspiral due to gravitational wave emission is adiabatic, we can also treat the change in magnetic field strength threading the black hole adiabatically, which makes the calculation surprisingly simple. Both of these assumptions will break down at some point prior to merger, so our main purpose will be to suggest this novel mechanism, present admittedly rough quantitative estimates, and motivate further investigation via numerical simulations. In particular, the dynamics of the magnetic field, which we are ignoring, is likely to become significant and interesting at some point prior to merger, and could have a substantial quantitative impact on the observed luminosities.

Throughout, we will use BH or NS to designate black hole or neutron star quantities, and M and r will refer to the mass and orbital separation of the system. We will focus on the late inspiral regime, where the gravitational waves are detectable and the magnetic field threading the black hole will be significant. This regime will be well within the neutron-star light cylinder, defined by $r_c \leq c/\Omega_{\rm NS}$, where $\Omega_{\rm NS}$ is the spin frequency of the neutron star, as well as the Alfvén radius, where the magnetic field dominates the plasma dynamics. Therefore, the black hole orbits within the neutron star magnetosphere, where a tenuous plasma and the magnetic field lines that thread it co-rotate with the neutron star [15]. We focus on typical, slowly spinning neutron stars with spin periods of $\mathcal{O}(1 \text{ sec})$, so that the co-rotating magnetic field is well-approximated by a dipole rotating at the neutron-star spin frequency. Assuming that the plasma in the neutron star magnetosphere is low-density and the force-free approximation applies [3, 15], and the magnetic field strength within the orbital plane is given by $|\vec{B}| = B_{\rm o}(r_{\rm NS}/r)^3$ corresponding to a typical dipole (where $r_{\rm NS}$ and $B_{\rm o}$ are the radius and surface magnetic field strength of the neutron star), there will be a substantial electromotive force (emf) between the horizon of the black hole and the poles of the neutron star. This emf will drive a current along magnetic field lines, creating an astronomical circuit, where the black hole acts like a battery that absorbs power due to its intrinsic resistivity and drives electromagnetic power toward the neutron star.

To estimate the power generated by the black hole-neutron star circuit, we begin with Faraday's law around a black hole:

$$\oint \alpha d\vec{l} \cdot \vec{E} = -\frac{d}{dt}\int \vec{B} \cdot d\vec{A} - \oint \alpha d\vec{l} \cdot (\vec{v} \times \vec{B}) \tag{1}$$

where α accounts for the gravitational redshift of fiducial observers ($\alpha \equiv \sqrt{1 - 2M_{\rm BH}/r}$ for a nonspinning black hole, where $M_{\rm BH}$ is the black-hole mass) and t is the time as measured by an observeer at rest at infinity. Of course, the binary does not yield a pure single black-hole spacetime but rather a spacetime better approximated by a post-Newtonian expansion with two bodies as sources. However, for our estimate we proceed with relativistic values for a single black hole to gauge the magnitude of the effect.

Faraday's law has the usual interpretation. An emf along a closed loop is related to the magnetic flux through an area enclosed by the loop and to the motion of the circuit relative to the magnetic field. While the law applies to any hypothetical loop we might draw, we are interested in the emf along the magnetic field lines, since they act as the wires for conduction of the plasma. The contour integration for employing Faraday's law in Eq. (1) spans one hemisphere of the black-hole horizon, and connects the black hole's equator with one magnetic pole of the neutron star by running parallel to the surrounding magnetic field lines (see Fig. 1). Assuming the rate of change of magnetic flux is negligible for this contour on orbital timescales, the emf is due solely to the relative motion of the contour through the neutron-star magnetic field. This contour yields no contribution from the segments that run parallel to the magnetic field, so that only the contribution from near the black-hole horizon remains. The emf, or equivalently the potential difference across either black-hole hemisphere, is given by

$$\begin{aligned} V &= \oint \alpha d\vec{l} \cdot \vec{E} = - \oint \alpha d\vec{l} \cdot (\vec{v} \times \vec{B}) \\ &= -2 r_{\rm H} B_{\rm o} \left(\frac{r_{\rm NS}}{r}\right)^3 \left[r\left(\Omega_{\rm orb} - \Omega_{\rm NS}\right) - \frac{a}{4\sqrt{2}} \right] \end{aligned} \tag{2}$$

where $a \equiv S_{\rm BH}/M_{\rm BH}^2$ and $r_{\rm H}$ are the dimensionless spin parameter and horizon radius of the black hole ($0 \leq a \leq 1$, and $r_{\rm H} = 2M$ for a nonspinning black hole) $v = [r(\Omega_{\rm orb} - \Omega_{\rm NS}) - a/(4\sqrt{2})]/\alpha$ is the azimuthal velocity of magnetic field lines as measured by non-inertial observers at rest with respect to the black hole (i. e. the fiducial observers, or "FIDOs", of the membrane paradigm [14]). $\Omega_{\rm orb}$ and $\Omega_{\rm NS}$ are the angular frequencies of the orbit and the neutron star spin, respectively, with $\Omega_{\rm NS}$ defined here to be positive for spin aligned with the orbital angular momentum, and negative for anti-aligned spins. Since the timescale for tidal locking greatly exceeds the inspiral timescale [16], realistic systems will not be tidally locked, and v will be generically non-zero even for nonspinning black holes. The term $a/(4\sqrt{2})$ is an approximation for the velocity of FIDOs as viewed from infinity due to the black-hole

spin, and bears further explanation. For a spinning black hole, the angular velocity of FIDOs depends on polar angle because the horizon is no longer spherical. However, for our purposes it is adequate to multiply the rms of the horizon radius, $r_{\rm H}/\sqrt{2}$, by the angular frequency relative to FIDOs of field lines around a single Kerr black hole, $\Omega_{\rm BH} = a/(4r_{\rm H})$, so that we will recover the standard Blandford-Znajek effect for $\Omega_{\rm orb} = \Omega_{\rm NS} = 0$.

The emf drives a current provided by the plasma so that the circuit obeys $V - IR_{\rm H} - IR_{\rm NS} = 0$ giving $I = V/(R_{\rm H} + R_{\rm NS})$ where $R_{\rm H}$ is the resistivity of the black hole membrane and $R_{\rm NS}$ is the resistivity of the neutron star. We use the remarkable result that the horizon has an effective resistance of 377 Ω (or 4π in geometrized units) [17]. As $R_{\rm H}$ is simply the resistance of free space, we are assuming throughout that the resistance due to the magnetosphere is negligible, and the resistance resulting from passing through the neutron star crust is comparable to $R_{\rm H}$, so that equal amounts of power are absorbed by the black hole through Ohmic dissipation as is emitted from near the horizon via a Poynting flux. Clearly, this assumption depends on the composition of the crust and the dynamics within the crust, which are very poorly constrained. We have also assumed that the black hole horizon is sufficiently small and far away that the dipole field is approximately vertical as it threads the black hole. This approximation will break down for black holes much more massive than the neutron star, or at separations near merger, but we are not interested in systems with very disparate masses ($M_{\rm BH} \gtrsim 500 {\rm M}_\odot$) since they will not be observable through gravitational waves with Advanced LIGO/Virgo and their electromagnetic luminosity will be mitigated by having a large black-hole mass. Since most of our other approximations will also break down close to merger, this should be a reasonable assumption.

We can now solve for the power generated by both hemispheres of the black-hole horizon:

$$L = 2\frac{V^2}{R_{\rm H}} = \frac{8\epsilon}{\pi}(\alpha v)^2 B^2 M_{\rm BH}^2 \tag{3}$$

$$= 3\times 10^{46}\epsilon \left(\alpha\frac{v}{c}\right)^2 \left(\frac{B_{\rm o}}{10^{12}\,{\rm G}}\right)^2 \left(\frac{r_{\rm NS}}{r}\right)^6 \left(\frac{M_{\rm BH}}{10\,{\rm M}_\odot}\right)^2 \ {\rm erg/s}\,. \tag{4}$$

We have assumed $\Omega_{\rm NS} \ll \Omega_{\rm orb}$, and we use $\epsilon \equiv (r_{\rm H}/2M)^2$ to account for the effect of black hole spin on the horizon size, so that ϵ ranges from unity for a nonspinning black hole to 1/4 for maximal spin. We choose 10 ${\rm M}_\odot$ as our fiducial black-hole mass because, in the absence of any observational constraints for black hole-neutron star binaries, we are left with results from population synthesis, which predict total mass distributions for these systems which peak at $\sim 8\,{\rm M}_\odot$ [18], and measurements from low mass X-ray binaries, which show narrow

distributions of $7.8 \pm 1.2\,\mathrm{M}_\odot$ [19]. The choice of a 10 $\mathrm{M}_\odot$ black hole paired with the standard choice of a 1.4 $\mathrm{M}_\odot$ neutron star then has the virtue of being a conservative estimate, since $L \propto M_{\rm BH}^2 r^{-6} \propto\!\!\!\sim M_{\rm BH}^{-4}$.

Before proceeding, we must point out a few more caveats for applying Eq. (4). We note that, if the neutron star spin period is $\mathcal{O}(\mathrm{ms})$, rather than the far more common periods of $\mathcal{O}(\mathrm{s})$, or if the separation is large, then the assumption that $\Omega_{\rm NS} \ll \Omega_{\rm orb}$ should not be employed, and the correct spin period should be used in Eq. (2). As millisecond pulsars generally have much weaker magnetic field strengths $B_{\rm o} \approx 10^9$ G, these systems are not only rarer, but also far weaker than the systems of primary interest in this work. In addition, for small stellar-mass black holes, the neutron-star radius or the tidal-disruption radius $r_{\rm tidal} \sim (M_{\rm BH}/M_{\rm NS})^{1/3} r_{\rm NS}$, rather than the light-ring radius ($r = 3M$ for a nonspinning black hole and $r = M$ for maximal spin), may set the minimum distance for applicability of Eq. (4), as the calculation plainly will break down before the orbital separation reaches these distances. Finally, it must be emphasized that the energy source for the luminosity predicted by Eq. (4) is a combination of the black hole's spin energy and the gravitational binding energy of the binary. Therefore, we can expect no clear relationship between Eq. (4) and the spin-down rate of the black hole or neutron star in this scenario, as L is non-vanishing (and, in fact, remains enormous) even in the limit that neither binary component is spinning. Furthermore, as the luminosity given by Eq. (4) is many orders of magnitude smaller than the gravitational wave luminosity for reasonable magnetic field strengths, the electromagnetic luminosity will not impact the dynamics. In Fig. 2, we show a comparison of the gravitational wave luminosity with the electromagnetic luminosity for both nonspinning and highly spinning cases.

While Eq. (4) will become increasingly inaccurate as the system approaches merger and magnetic field dynamics likely become significant, Eq. (4) could yield a reasonable approximation at separations outside the light ring if we use an accurate trajectory informed by numerical relativity simulations. For separations less than the radius of the light ring, a significant fraction of the Poynting flux predicted by Eq. (4) may be unable to escape from the black hole, so that the peak luminosity is likely to occur at a separation near the light ring. Therefore, in Fig. 2, we calculate the electromagnetic luminosity given by Eq. (4), with $v(t)$ calculated by evolving the pseudo-4PN Hamiltonian introduced in [20] for the case of $a = 0$, and by evolving the Hamiltonian introduced in [21] for the case of $a = 0.9$. The

spinning black hole in our example has a mass $M_{\rm BH} = 14\,{\rm M}_\odot$, because this specific set of system parameters is not expected to result in the tidal disruption of the neutron star [22]. Since it is near the disruption threshold, it should be seen as a near-optimal case, although it may nonetheless be more typical given the expected distribution of black hole spins. We include the gravitational wave energy flux to third post-Newtonian order in the Taylor expansion [23]. We truncate the electromagnetic luminosity of both systems at their respective light rings, where we estimate the light ring location by finding the radius of a particle on a circular orbit with infinite momentum. For comparison, we also calculate the gravitational wave luminosity for the nonspinning case using the same inspiral model combined with the implicit-rotating-source (IRS) merger model first presented in [24]. As all of these models were tuned to accurately match numerical relativity simulations of black hole-black hole binaries, they do not include any finite size effects from the neutron star. Although the models will therefore accumulate small differences from an actual simulation of the systems of interest, these errors will be far smaller at late times than the other approximations we have described in deriving Eq. (4). Naturally, for sufficiently small black-hole masses and sufficiently large black-hole spins, the neutron star will be tidally disrupted prior to merger, so our approximations will fail sometime before the system reaches the disruption radius. However, nonspinning black holes with $M_{\rm BH} = 3\,{\rm M}_\odot$ will not disrupt the neutron star prior to merger for realistic equations of state, and near-extremal spinning black holes will not disrupt unless $M_{\rm BH} < 10\,M_{\rm NS}$ (Fig. 4 of [22]), so the process we describe may be much more common than the signature from tidal disruption and subsequent accretion of the neutron star by the black hole.

We can use the results shown in Fig. 2 to estimate the peak luminosity for each system. Using the light-ring values for each case, we find

$$L_{\rm peak}(a=0) \approx 2\times 10^{42} \left(\frac{B_{\rm o}}{10^{12}\,{\rm G}}\right)^{2} \left(\frac{M_{\rm BH}}{10\,{\rm M}_\odot}\right)^{-4} \ {\rm erg/s} \tag{5}$$

for the nonspinning case, and

$$L_{\rm peak}(a=0.9) \approx 9\times 10^{43} \left(\frac{B_{\rm o}}{10^{12}\,{\rm G}}\right)^{2} \left(\frac{M_{\rm BH}}{10\,{\rm M}_\odot}\right)^{-4} \ {\rm erg/s} \tag{6}$$

for the spinning case. It is encouraging to note that Eq. (6) is approaching the Blandford-Znajek luminosity of a single maximally spinning black hole with the same total mass M, immersed in a magnetic field with the same strength $B_{\rm o}$ as our fiducial surface field. This is

consistent with our expectation that the luminosity as $r \to M$ and $v_{\rm orb} \to c$ for the binary should equal the luminosity for a maximally- spinning black hole, keeping all other variables equal, so that this observation is a useful sanity check.

While these estimates do not include any contribution from the final merger, it is still instructive to integrate Eq. (4) over the domain $r \geq r_{\rm LR}$, where $r_{\rm LR}$ is the light-ring radius, as a conservative estimate of the total energy emitted strictly by the process described herein, which yields

$$E_{\rm tot}(a=0) = \int_{r=r_{\rm LR}}^{\infty} dr\, \frac{L(r)}{dr/dt} \approx 3 \times 10^{39} \left(\frac{B_{\rm o}}{10^{12}\,{\rm G}}\right)^2 \left(\frac{M_{\rm BH}}{10\,{\rm M}_\odot}\right)^{-3} \text{erg} \tag{7}$$

for the nonspinning case, and

$$E_{\rm tot}(a=0.9) \approx 8 \times 10^{41} \left(\frac{B_{\rm o}}{10^{12}\,{\rm G}}\right)^2 \left(\frac{M_{\rm BH}}{10\,{\rm M}_\odot}\right)^{-3} \text{erg} \tag{8}$$

for the spinning case, where we again emphasize that this calculation is intended as an order-of-magnitude estimate, though it has the virtue of almost certainly being an underestimate, as it neglects the merger.

In the case of the Blandford-Znajek prototype, the system model would now be complete. However, in the scenario presented here, the luminosity given by Eq. (4) is not simply emitted as a jet, as it would be for Blandford-Znajek or for orbiting black-hole binaries with a uniform magnetic field geometry. Because the plasma in the magnetosphere is constrained to follow magnetic field lines, it will be channeled from just outside the black-hole horizon to the neutron-star magnetic poles. As the plasma is accelerated forward by the Poynting flux and transversely by the curving field lines of the dipole geometry, some fraction (which we will label η) of the power will be dissipated as synchrotron and curvature radiation that is strongly beamed by relativistic effects (see Fig. 1). Because the intensity of the curvature radiation depends only on the magnetic field geometry, and the intensity of the synchrotron radiation depends on the conserved magnetic flux, which determines the fraction of electrons and positrons populating excited Landau levels, we expect the intensity of the emitted radiation as the plasma transits from the black hole to the neutron star surface will be fairly uniform. The synchrotron and curvature radiation will therefore uniformly sweep out the entire plane that contains the black hole and neutron star and that is parallel to the magnetic field. As this plane rotates with the orbiting binary, we have the interesting phenomena of a beamed jet which covers the entire 4π sky each orbit. Since the orbital

timescale and the burst timescale (defined as the time interval containing 90% of the total energy) are both $\mathcal{O}$(ms), the entire sky should receive a brief blast of luminosity at the magnitude of Eqs. (5) and (6). Judging from the behavior observed in numerical simulations of neutron star magnetospheres accelerating a driven plasma wind [25], the majority of the observed luminosity due to acceleration of the plasma within the magnetosphere will likely be due to curvature radiation, which will dominate above 10 MeV, with synchrotron radiation accounting for the majority of observed radiation in hard X-rays and soft γ-rays, with energies 10 keV $\lesssim h\nu \lesssim$ 10 MeV.

As some unknown fraction of the emitted energy will still be in the form of plasma kinetic energy when the plasma reaches the neutron star, the flux of plasma may induce a hot spot on the neutron-star surface (see Fig. 1). Assuming even a small fraction of the luminosity predicted by Eq. (4) is so-deposited, the thermal blackbody emission re-radiated by the neutron star surface will be substantially super-Eddington. Assuming a hot spot surface area $A \approx 1$ km^2, the temperature of the hot spot will be given by

$$T = \left(\frac{L_{\rm BB}}{\sigma A}\right)^{1/4} = 1\,{\rm MeV}\,[\epsilon(1-\eta)]^{1/4}\sqrt{\left(\alpha\frac{v}{c}\right)\left(\frac{B_{\rm o}}{10^{12}\,{\rm G}}\right)\left(\frac{r_{\rm NS}}{r}\right)^3\left(\frac{M}{10\,{\rm M}_\odot}\right)}, \tag{9}$$

where $L_{\rm BB} \equiv (1-\eta)L$ is the fraction of total luminosity (as given by Eq. (4)) that is emitted from the neutron surface and σ is the Stefan-Boltzmann constant. The maximum temperature of the hotspot, which results from inserting the maximum luminosities from Eqs. (5) and (6) into Eq. (9), is given by

$$T_{\rm max}(a=0) = 100\,{\rm keV}\,(1-\eta)^{1/4}\sqrt{\frac{B_{\rm o}}{10^{12}\,{\rm G}}\left(\frac{M}{10\,{\rm M}_\odot}\right)^{-1}}\,, \tag{10}$$

for the nonspinning case, and

$$T_{\rm max}(a=0.9) = 300\,{\rm keV}\,(1-\eta)^{1/4}\sqrt{\frac{B_{\rm o}}{10^{12}\,{\rm G}}\left(\frac{M}{10\,{\rm M}_\odot}\right)^{-1}}\,, \tag{11}$$

for the spinning case. Therefore, the total luminosity observed from this process will occur entirely within the hard X-ray–soft γ-ray bands. This emission should therefore be observable by the Swift γ-ray burst and Fermi γ-ray space telescopes. It is noteworthy that there is some evidence for a subclass of short GRBs with durations less than 100 ms, photon energies near the upper bound of short GRBs, and no apparent afterglow [26], which is quite consistent with our expectations for the process we have described, and which is quite distinct

from the expectations from neutron star-neutron star mergers and the tidal disruption and subsequent accretion of neutron stars by low-mass black holes. More speculatively, we note that these systems could bear some resemblance to the behavior observed in the subset of soft γ-ray repeaters (SGRs) that display a regular sequence of pulses occurring at $\mathcal{O}(1\text{ sec})$ intervals. The leading candidate for driving SGRs is the onset of a star quake on a magnetar [27]. It is interesting to note, however, that if we take the observed SGR periodicity to be the orbital period, and assume $B_{\rm o} \approx 10^{15}$ G which is consistent with magnetars, then Eq. (4) predicts a luminosity of $\mathcal{O}(10^{35}$ erg/s, which agrees very well with the observed luminosity from this subset of SGRs in quiescence [28]. Clearly, this could simply be an interesting coincidence.

Given the brevity of the signal, the Swift Burst Alert Telescope (BAT) [29] or the Fermi GLAST Burst Monitor (GBM) [30] would be the most suitable instruments for observing these events. The 100 μs and 2 μs timing resolutions of BAT and GBM, respectively, are more than adequate to resolve a signal with a duration of $\mathcal{O}(\text{ms})$, and their large fields-of-view are critical for observing these transients. We can use the flux sensitivities ($F_{\rm min} = 0.07$ photons/cm^2/s for BAT and 0.7 photons/cm^2/s for GBM) to estimate the detectable range for these events. Using BAT and Eq. (5) and assuming the event is beamed into a cone covering a solid angle $\Delta\Omega = 100\,\text{deg}^2$, the maximum luminosity distance range is given by

$$D_L = \sqrt{\frac{L_{\rm peak}}{\Delta\Omega F_{\rm min}}} = 60\sqrt{\frac{100\,\text{deg}^2}{\Delta\Omega}}\,\text{Mpc}\,. \tag{12}$$

If we instead assume a narrower beaming $\Delta\Omega = 1\,\text{deg}^2$, or assume a neutron-star surface magnetic field of $\mathcal{O}(10^{14}\text{ G})$, the range would be 600 Mpc, or a redshift $z = 0.1$ using WMAP seven-year cosmological parameters [31], which is comparable to the upper distance limit for gravitational wave detection of these systems with Advanced LIGO (see Fig. 4b of [32]). As the estimates given by Eqs. (5) and (6) are likely to be exceedingly conservative estimates of the true peak luminosity for this process (which will likely occur closer to merger), this justifies our claim that the process we describe could potentially drive an electromagnetic counterpart to any gravitational wave signal from a black hole-neutron star binary that is observable by Advanced LIGO/Virgo.

Despite the uncertainties in the expected signature of these sources, we would expect that the electromagnetic emission would consist of a synchrotron/curvature hard X-ray/soft γ-ray component whose intensity varies over the orbital period, and a blackbody component

with a peak energy corresponding to soft γ-rays, whose intensity varies over an interval aprroximately given by the sum of the orbital period, neutron star spin period, and black hole spin period. The total electromagnetic signature could therefore have two distinct observable components rising and falling with different periods if the source is close enough. Both the synchrotron/curvature and the blackbody emission will be relativistically beamed, which will increase the potential range for observing these events, but may decrease the event rate for observations within a fixed volume. If the synchrotron/curvature component dominates, the beaming may not decrease the event rate, given that the unusual geometry of the system will blanket the entire sky with radiation over intervals comparable to the burst duration. The numerical simulation of these systems could be done with existing codes, and would be an invaluable confirmation (or refutation) of the mechanism which we have described.

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Acknowledgements We wish to thank Cole Miller, Frans Pretorius, and Anatoly Spitkovsky for helpful feedback on the manuscript. This work was supported by NSF grant AST-0908365. JL gratefully acknowledges a KITP Scholarship.

Competing Interests The authors declare that they have no competing financial interests.

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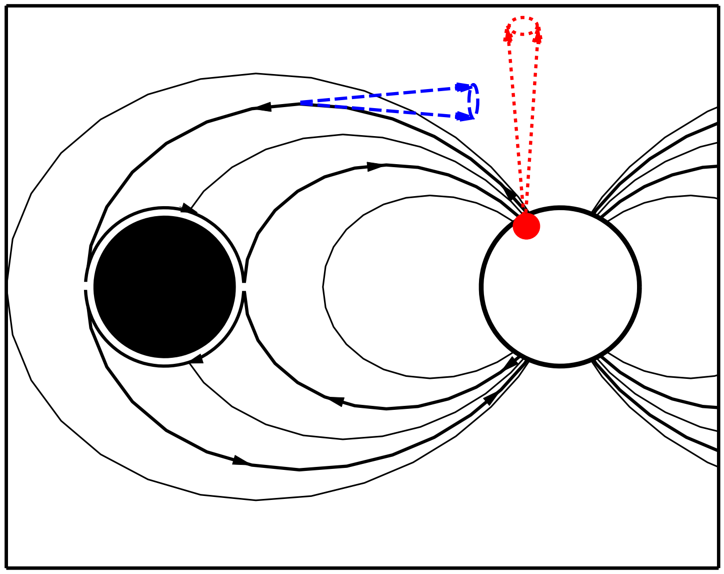

FIG. 1: A black hole-neutron star binary, where the black hole (black) is threaded by the dipole magnetic field of the neutron star (white). The thick lines with arrows show the integration contours used in Eq. (2), including the segment that runs over the horizon, just outside the black hole [14]. The Poynting flux from the black hole imparts kinetic energy into the plasma, which in turn emits synchrotron radiation, and the magnetic field geometry accelerates the plasma and generates curvature radiation. The synchrotron/curvature radiation will be relativistically beamed (blue dashed cone). The plasma that reaches the neutron star surface will induce radiation as its remaining kinetic energy dissipates by heating the surface. This radiation will be nearly blackbody, and will likely be relativistically beamed as well (red dotted cone).

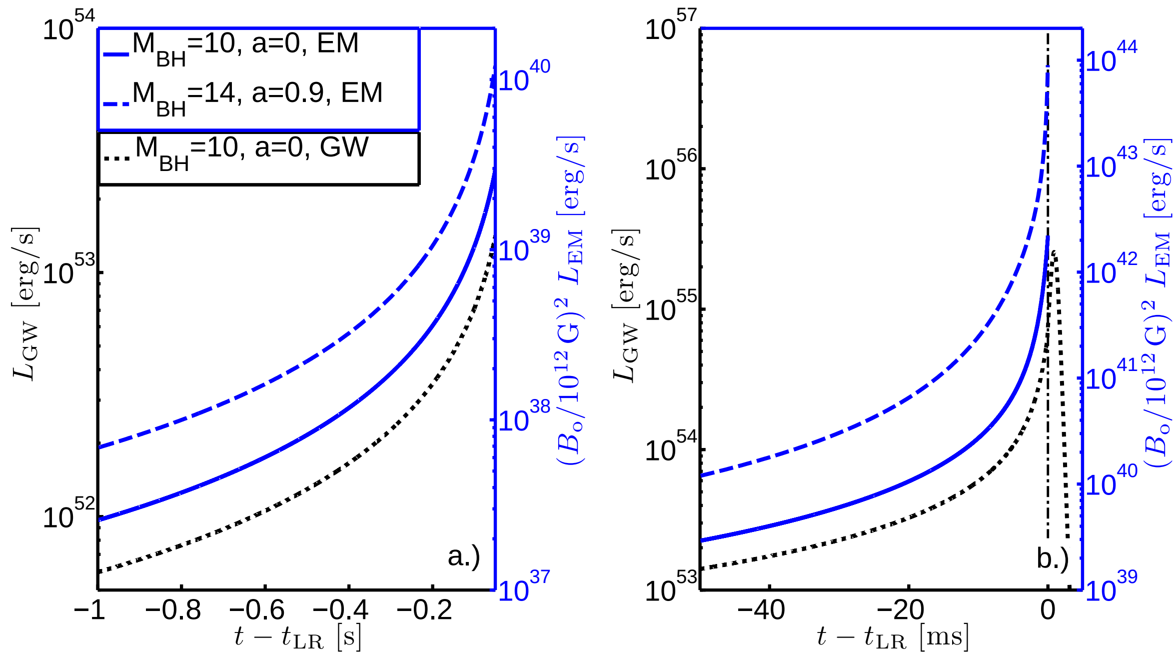

FIG. 2: A comparison of the electromagnetic luminosity for our fiducial case (nonspinning black hole, $M_{\mathrm{BH}} = 10\,\mathrm{M}_{\odot}$) and our most luminous case ($a \equiv S/M_{\mathrm{BH}}^2 = 0.9$, $M_{\mathrm{BH}} = 14\,\mathrm{M}_{\odot}$) with the gravitational wave luminosity for our fiducial case. In (a.), we show the second preceding the final burst, which is the typical timescale for short GRBs. In (b.), we show the final 50 ms of the burst, which dominates the total energy output. We note that the gravitational wave luminosity rises considerably after the light ring (designated by a thin vertical dash-dotted line), so it seems likely that the peak electromagnetic emission from the process we describe may substantially excede that predicted by Eqs. (5) and (6), though numerical simulation will be required to explore this regime.